

Constraints on dark matter annihilation to fermions and a photon

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We consider Majorana dark matter annihilation to fermion - anti-fermion pair and a photon in the effective field theory paradigm, by introducing dimension 6 and dimension 8 operators in the Lagrangian. For a given value of the cut-off scale, the latter dominates the annihilation process for heavier dark matter masses. We find a cancellation in the dark matter annihilation to a fermion - anti-fermion pair when considering the interference of the dimension 6 and the dimension 8 operators. Constraints on the effective scale cut-off is derived while considering indirect detection experiments and the relic density requirements and then comparing them to the bound coming from collider experiments.

I. INTRODUCTION

Astrophysical observations at all scales confirm the presence of dark matter [1]. The ubiquitous astrophysical discovery of dark matter from the solar system scale to the cosmic microwave background scale confronts us with many profound mysteries of nature [2–6]. Although astrophysical observations give us important clues on dark matter properties, precise questions about dark matter can only be answered once we determine the particle properties of dark matter.

Numerous dark matter particle candidates exist with masses ranging from $\sim 10^{-22}$ eV to $10^{-8} M_\odot$, but perhaps the most widely searched for particle goes under the generic name of weakly interacting massive particles. Weakly interacting massive particles have masses in between a few GeV to few tens of TeV, and interact with the Standard Model particles with “weak” strength [7–9]. These properties make them perfect candidates for searches in colliders [10], indirect detection [11–21], and direct detection [22–31].

In indirect detection of dark matter, we search for faint signals of dark matter annihilation and decay from the cosmos amidst the overwhelming conventional astrophysical background. Indirect detection of dark matter can give us information about dark matter properties which are not easily accessible otherwise [20]. Due to the enormous and varied astrophysical background, it is often necessary to either search for a signal unique to the dark matter particle candidate or devise clever search strategies [32].

Dark matter interactions to the Standard Model

sector can be parametrized by higher dimensional non-renormalizable operators using the effective field theory technique [33–44]. These operators are suppressed by different powers of the effective cut-off scale Λ . If the new physics is sufficiently decoupled from the Standard Model particles then this is an adequate description. There have been recent discussion about the validity of effective field theory, especially at the colliders [45–48]. In spite of its limitations, effective field theory approach to dark matter is useful to classify various interactions and make progress in understanding dark matter physics.

Various well motivated new physics extension of the Standard Model of particle physics predict that the dark matter particle is a Majorana fermion. General considerations suggest that the annihilation of Majorana dark matter particles to fermions in the s -wave channel is proportional to $(m_f/m_\chi)^2$, where m_f and m_χ stands for the mass of the Standard Model fermion and dark matter particle respectively [49–51]. The rate in the p -wave channel is proportional to v^2 , where v is the relative velocity of two incoming dark matter particles.

Given that the dark matter particle is expected to have a mass $\gtrsim 100$ GeV, and almost all the Standard Model fermions (except the top quark) have a mass of < 5 GeV, this implies a huge suppression of the p -wave annihilation rate. The dark matter velocity in the Solar circle is $\sim 10^{-3}c$, and the typical velocities in clusters and dwarf galaxies are $\sim 10^{-2}c$ and $\sim 10^{-4}c$ respectively. These small velocities ensure that the p -wave contribution to non-relativistic Majorana dark matter particles annihilating to fermions is velocity suppressed.

A possible way out of this conundrum was suggested long ago in which a photon is radiated out from the final state fermion or the charged mediator and this lifts the suppression [52, 53]. This contribution has also been included in numerical packages which calculate dark matter properties for supersymmetric dark matter

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candidates [54]. Recently it has also been realized that in general the radiation of electroweak gauge bosons lift this suppression [55–67]. Additional work has also been done regarding gluon [68, 69] and Higgs radiation [70].

It is tempting to ask if it is possible to incorporate boson bremsstrahlung in the framework of effective field theory of dark matter. Ref. [41] presented a list of operators which describes Majorana dark matter annihilation to fermion - anti-fermion pair and an electroweak gauge boson. It was shown that dimension-8 operator is required to consistently describe a Majorana dark matter annihilation into a fermion - anti-fermion pair and an electroweak gauge boson.

In this present work, using these operators, constraints on Λ obtained from the relic density requirement and present limits on dark matter annihilation is compared with those obtained from collider. The operators that we consider do not have a non-relativistic limit and hence the constraints from dark matter direct detection experiments do not apply [40].

In Section II we discuss the effective operators and compute the cutoff scale Λ to obtain the correct dark matter relic density. We compare these values with the constraints from colliders and indirect detection experiments. We conclude in Section III.

II. EFFECTIVE OPERATOR MODEL

We assume that the dark matter particle is a Majorana fermion and is represented by χ . The lowest order interaction between dark matter particles and Standard Model fermions, denoted by f , that we consider is of the form [41]

$$\mathcal{L}_{d=6} = \frac{1}{\Lambda^2} (\bar{\chi} \gamma^5 \gamma^\mu \chi) (\bar{f} \gamma_\mu f). \quad (1)$$

The higher order Lagrangian that we consider is of the form [41]

$$\begin{aligned} \mathcal{L}_{d=8} = & \frac{1}{\Lambda^4} (\bar{\chi} \gamma^5 \gamma^\mu \chi) \left[\left(\bar{f}_L \overleftarrow{D}_\rho \right) \gamma_\mu \left(\overrightarrow{D}^\rho f_L \right) \right. \\ & \left. + \left(\bar{f}_R \overleftarrow{D}_\rho \right) \gamma_\mu \left(\overrightarrow{D}^\rho f_R \right) \right], \end{aligned} \quad (2)$$

where following Ref. [41], we define

$$\bar{f}_L \overleftarrow{D}_\mu = (\partial_\mu \bar{f}_L) - ig \frac{\sigma^i}{2} W_\mu^i \bar{f}_L - ig' Y_f B_\mu \bar{f}_L, \quad (3)$$

$$\overrightarrow{D}_\mu f_L = (\partial_\mu f_L) + ig \frac{\sigma^i}{2} W_\mu^i f_L + ig' Y_f B_\mu f_L, \quad (4)$$

$$\bar{f}_R \overleftarrow{D}_\mu = (\partial_\mu \bar{f}_R) - ig' Y_f B_\mu \bar{f}_R, \quad (5)$$

$$\overrightarrow{D}_\mu f_R = (\partial_\mu f_R) + ig' Y_f B_\mu f_R, \quad (6)$$

$$(7)$$

where σ^i denotes the Pauli matrices, W^μ and B^μ denote the Standard Model gauge bosons, Y_f denotes

the hypercharge and g and g' denote the Standard Model SU(2) and U(1) gauge couplings respectively.

The total Lagrangian is given by

$$\mathcal{L} = d_6 \mathcal{L}_{d=6} + d_8 \mathcal{L}_{d=8}, \quad (8)$$

where d_6 and d_8 are arbitrary complex constants.

A. $\chi + \chi \rightarrow f + \bar{f}$

We will first calculate the cross section for the process $\chi(k_1) + \chi(k_2) \rightarrow f(p_1) + \bar{f}(p_2)$, where we denote the four momentum of the particle in parenthesis. An important ingredient in our calculation is that we take into account both the dimension 6 and dimension 8 terms. As we will see due to the interference of these two terms, there is a cancellation feature in the σv which is not present when one only considers dimension 6 term. From Eq. (8), the relevant part of the Lagrangian responsible for the process is

$$\begin{aligned} \mathcal{L} = & \frac{d_6}{\Lambda^2} (\bar{\chi} \gamma^5 \gamma^\mu \chi) (\bar{f} \gamma_\mu f) \\ & + \frac{d_8}{\Lambda^4} (\bar{\chi} \gamma^5 \gamma^\nu \chi) \left[(\partial_\rho \bar{f}) \gamma_\nu (\partial^\rho f) \right]. \end{aligned} \quad (9)$$

From this Lagrangian, we obtain the following cross section for the process $\chi(k_1) + \chi(k_2) \rightarrow f(p_1) + \bar{f}(p_2)$

$$\begin{aligned} \sigma = & \frac{1}{v} \frac{v^2}{3\pi\Lambda^8} \sqrt{1 - \frac{m_f^2}{m_\chi^2}} (2m_\chi^2 + m_f^2) \left\{ |d_6|^2 \Lambda^4 + |d_8|^2 (2m_\chi^2 + m_f^2)^2 \right. \\ & \left. - \Lambda^2 (d_6 d_8^* + d_6^* d_8) (2m_\chi^2 - m_f^2) \right\}, \end{aligned} \quad (10)$$

where m_χ and m_f denote the mass of the dark matter particle and Standard Model fermion respectively. From Eq. (10), one can see that due to the v^2 dependence of σv , constraints from dark matter annihilation in the present epoch will be very weak. Having an additional vector boson in the final state removes this v^2 dependence at leading order [41, 56–58, 60].

B. $\chi + \chi \rightarrow f + \bar{f} + \gamma$

We now calculate the cross-section for the process $\chi(k_1) + \chi(k_2) \rightarrow f(p_1) + \bar{f}(p_2) + \gamma(k)$. For completeness we will show the explicit steps in our calculation. From Eq. (2), the connection between the dark matter particle and the vector bosons come from the covariant derivative: $D_\mu \equiv \partial_\mu - ig W_\mu^i (\sigma^i/2) - ig' Y_f B_\mu$. This can also be written as $D_\mu \equiv \partial_\mu - i(g/\sqrt{2})(W_\mu^+ T^+ + W_\mu^- T^-) - i(g/\cos\theta_W)Z_\mu(T_3 - \sin^2\theta_W Q) - ieQ A_\mu$. To derive this, we use $W_\mu^\pm = (1/\sqrt{2})(W_\mu^1 \mp iW_\mu^2)$, $T^i = \sigma^i/2$, $T^\pm = T^1 \pm iT^2$, $Z_\mu^0 = \cos\theta_W W_\mu^3 - \sin\theta_W B_\mu$, $A_\mu = \sin\theta_W W_\mu^3 +$

$\cos\theta_W B_\mu$, and $\cos\theta_W = g/(\sqrt{g^2 + g'^2})$. We denote the photon by A_μ and the weak mixing angle by θ_W .

The effective Lagrangian for the given process is

$$\begin{aligned}\mathcal{L}_{\text{eff}} = & \frac{d_6}{\Lambda^2} (\bar{\chi} \gamma^5 \gamma^\mu \chi) (\bar{f} \gamma_\mu f) \\ & + \frac{d_8}{\Lambda^4} (\bar{\chi} \gamma^5 \gamma^\nu \chi) \left[(\partial_\rho \bar{f}) \gamma_\nu (\partial^\rho f) \right. \\ & \left. + i e Q A_\rho \left\{ (\partial^\rho \bar{f}) \gamma_\nu f - \bar{f} \gamma_\nu (\partial^\rho f) \right\} \right]. \quad (11)\end{aligned}$$

$$\begin{aligned}\mathcal{M} = & -i e Q \bar{u}(k_1) \gamma^5 \gamma^\mu v(k_2) \left[\left(\frac{d_6}{\Lambda^2} - \frac{d_8}{\Lambda^4} p_2 \cdot (p_1 + k) \right) \bar{u}(p_1) \not{\epsilon}(k) \frac{p_1 + k + m_f}{(p_1 + k)^2 - m_f^2 + i\epsilon} \gamma_\mu v(p_2) \right. \\ & + \left(\frac{d_6}{\Lambda^2} - \frac{d_8}{\Lambda^4} p_1 \cdot (p_2 + k) \right) \bar{u}(p_1) \gamma_\mu \frac{-p_2 - k + m_f}{(p_2 + k)^2 - m_f^2 + i\epsilon} \not{\epsilon}(k) v(p_2) \\ & \left. - \frac{d_8}{\Lambda^4} \bar{u}_{p_1} \epsilon_\rho(k) \gamma_\mu v_{p_2} (p_1^\rho - p_2^\rho) \right]. \quad (12)\end{aligned}$$

The amplitude contains the emission of photon from the fermion due to the operator containing the coefficients d_6 and d_8 . It also includes the emission of photon from the blob due to the operator containing the coefficient d_8 . The square of the amplitude includes 9 terms and we will not present it here for brevity. The calculation of the cross section involves integration of the square of the amplitude over the 3-body phase space [71, 72].

The 3-body phase space can be written as

$$\begin{aligned}d(3PS) = & \frac{d^3 \mathbf{p}_1}{(2\pi)^3 2E_1} \frac{d^3 \mathbf{p}_2}{(2\pi)^3 2E_2} \frac{d^3 \mathbf{k}}{(2\pi)^3 2E_\gamma} \\ & \times (2\pi)^4 \delta^{(4)}(k_1 + k_2 - p_1 - p_2 - k), \quad (13)\end{aligned}$$

where E_1 , E_2 , and E_γ denote the energy component of the 4-momentum p_1 , p_2 , and k respectively. We simplify this phase space by decomposing it into a product of 2-body phase spaces. Let us denote $p_{1k} = p_1 + k$, and $p_{1k}^2 = m_{1k}^2$. We insert the identity $\int \frac{d^4(p_{1k})}{(2\pi)^4} (2\pi)^4 \delta(p_{1k} - p_1 - k) \Theta(p_{1k}^0) = 1$ and $\int \frac{d(m_{1k}^2)}{(2\pi)} (2\pi) \delta(p_{1k}^2 - m_{1k}^2) = 1$ in Eq. (13), where p_{1k}^0 is the 0th component of the 4-vector p_{1k} . The resulting expression can be simplified by noting that $\int \frac{d^4(p_{1k})}{(2\pi)^4} 2\pi \delta(p_{1k}^2 - m_{1k}^2) \Theta(p_{1k}^0) = \int \frac{d^3 \mathbf{p}_{1k}}{(2\pi)^3 2\sqrt{\mathbf{p}_{1k}^2 + m_{1k}^2}}$.

The relevant Feynman diagrams are shown in Fig. 1. The amplitude for the process is given by

The 3-body phase space can then be written as

$$\begin{aligned}d(3PS) = & \frac{d^3 \mathbf{p}_1}{(2\pi)^3 2E_1} \frac{d^3 \mathbf{p}_2}{(2\pi)^3 2E_2} \frac{d^3 \mathbf{k}}{(2\pi)^3 2E_\gamma} \frac{d^3 \mathbf{p}_{1k}}{(2\pi)^3 2\sqrt{p_{1k}^2}} \\ & \times \frac{d(p_{1k}^2)}{2\pi} (2\pi)^4 \delta^{(4)}(k_1 + k_2 - p_{1k} - p_2) \\ & \times (2\pi)^4 \delta^{(4)}(p_{1k} - p_1 - k), \quad (14)\end{aligned}$$

where $p_{1k}^2 = \mathbf{p}_{1k}^2 + m_{1k}^2$.

The three-dimensional integrals in Eq. (14) is Lorentz-invariant and can be calculated in any frame. A compact expression for these can be obtained if any two of the three-dimensional integrals are integrated in the center of momentum frame of the two momentum vectors involved. Following this strategy, we obtain the following expression in the center of momentum frame of p_{1k} :

$$\begin{aligned}& \frac{d^3 \mathbf{p}_1}{(2\pi)^3 2E_1} \frac{d^3 \mathbf{k}}{(2\pi)^3 2E_\gamma} (2\pi)^4 \delta^{(4)}(p_{1k} - p_1 - k) \\ = & \frac{d\phi_{p_{1k}} d(\cos \theta_{p_{1k}})}{16\pi^2} \frac{1}{2([p_1^0 + k^0]_{p_{1k}})^2} \left[([p_1^0 + k^0]_{p_{1k}})^4 + p_1^4 \right. \\ & + k^4 - 2p_1^2 k^2 - 2p_1^2 ([p_1^0 + k^0]_{p_{1k}})^2 \\ & \left. - 2k^2 ([p_1^0 + k^0]_{p_{1k}})^2 \right]^{1/2}, \quad (15)\end{aligned}$$

where $\phi_{p_{1k}}$ and $\theta_{p_{1k}}$ denote the spherical polar coordinates in the center of momentum frame of p_{1k} . The 0th component of the 4 vectors p_1 and k are denoted by p_1^0 and k^0 . In the center of momentum frame of p_{1k} , we have $\mathbf{p}_{1k} = 0$, so that $p_{1k}^2 = (p_{1k}^0)^2$. This will simplify the expression in Eq. (15) where we replace $([p_1^0 + k^0]_{p_{1k}})^2$ by p_{1k}^2 .

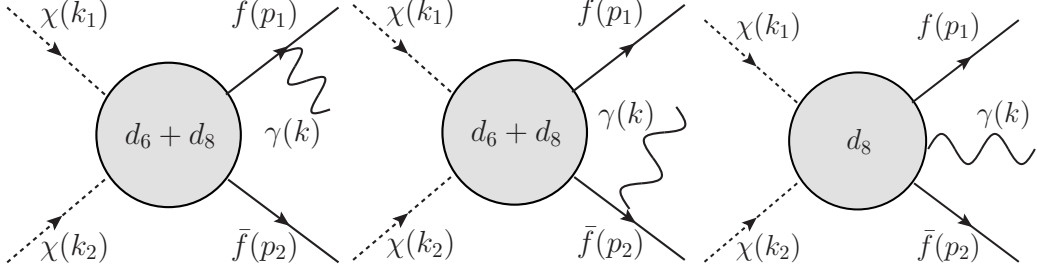


FIG. 1. Feynman diagrams for dark matter annihilation into a fermion - anti-fermion pair and a photon. The dark matter, fermion, anti-fermion and the photon are denoted by χ , f , \bar{f} and γ respectively. The four-momentum associated with each particle is given in parenthesis next to the particle. The operator which contributes to the Feynman diagram is written in the blob.

Similarly we can also write

$$\begin{aligned} & \frac{d^3 \mathbf{p}_2}{(2\pi)^3 2E_2} \frac{d^3(\mathbf{p}_{1k})}{(2\pi)^3 2\sqrt{\mathbf{p}_{1k}^2 + m_f^2}} (2\pi)^4 \delta^{(4)}(P - p_{1k} - p_2) \\ &= \frac{d\phi_P d(\cos\theta_P)}{32\pi^2 P^2} \left[P^4 + p_2^4 + (p_{1k})^4 - 2p_2^2(p_{1k})^2 \right. \\ & \quad \left. - 2p_2^2 P^2 - 2P^2(p_{1k})^2 \right]^{1/2}, \end{aligned} \quad (16)$$

where $P = p_2 + p_{1k} = k_1 + k_2$. The spherical polar coordinates in the center of momentum frame of P are denoted by ϕ_P and θ_P .

The minimum value of $p_{1k}^2 = (p_1 + k)^2$ is m_f^2 . We choose our coordinate system such that the integral is symmetric about $\phi_{p_{1k}}$. After integrating over $\phi_{p_{1k}}$ we obtain

$$\begin{aligned} d(3PS_\phi) &= \int_{p_{1k}^2=m_f^2}^s \left\{ s^2 + p_{1k}^4 + m_f^4 - 2s p_{1k}^2 - 2s m_f^2 \right. \\ & \quad \left. - 2m_f^2 p_{1k}^2 \right\}^{1/2} \left\{ p_{1k}^4 + m_f^4 + k^4 - 2m_f^2 p_{1k}^2 \right. \\ & \quad \left. - 2m_f^2 k^2 - 2p_{1k}^2 k^2 \right\}^{1/2} \\ & \quad \times \frac{1}{2^6} \frac{1}{s p_{1k}^2} d\phi d(\cos\theta_P) d(\cos\theta_{p_{1k}}) \frac{d(p_{1k}^2)}{(2\pi)^4}, \end{aligned} \quad (17)$$

where we now re-denote ϕ_P as ϕ .

We can now choose the individual 3-vectors as

$$\begin{aligned} \left[\mathbf{k}_1 \right]_P &= \frac{\sqrt{s}}{2} |\mathbf{v}_\chi| (\sin\theta_P \cos\phi, \sin\theta_P \sin\phi, \cos\theta_P), \\ \left[\mathbf{k}_2 \right]_P &= -\frac{\sqrt{s}}{2} |\mathbf{v}_\chi| (\sin\theta_P \cos\phi, \sin\theta_P \sin\phi, \cos\theta_P), \\ \left[\mathbf{p}_2 \right]_P &= (0, 0, -\sqrt{(E_2^P)^2 - m_f^2}), \\ \left[\mathbf{p}_{1k} \right]_P &= (0, 0, \sqrt{(E_2^P)^2 - m_f^2}), \end{aligned}$$

$$\begin{aligned} \left[\mathbf{p}_1 \right]_{p_{1k}} &= \sqrt{(E_1^{p_{1k}})^2 - m_f^2} (\sin\theta_q, 0, \cos\theta_q), \\ \left[\mathbf{k} \right]_{p_{1k}} &= -\sqrt{(E_1^{p_{1k}})^2 - m_f^2} (\sin\theta_q, 0, \cos\theta_q), \end{aligned} \quad (18)$$

where $[\dots]_P$ and $[\dots]_{p_{1k}}$ denote evaluation in the laboratory and p_{1k} rest frame respectively. The Mandelstam variable $s = (k_1 + k_2)^2 = (p_1 + p_2 + k)^2 = (p_{1k} + p_2)^2$, $t_1 = (k_1 - p_{1k})^2 = (k_2 - p_2)^2$.

By definition, $p_2^0 = E_2^P$, so that $k_2 \cdot p_2 = m_\chi E_2^P - m_\chi |v_\chi| \cos\theta_P \sqrt{(E_2^P)^2 - m_f^2}$. Since $2p_2 \cdot p_{1k} = s - m_f^2 - p_{1k}^2 = 2E_2^P p_{1k}^0 - 2 \left(-\sqrt{(E_2^P)^2 - m_f^2} \right) \left(\sqrt{(E_2^P)^2 - m_f^2} \right)$, and $p_{1k}^0 = P^0 - p_2^0 = 2m_\chi - E_2^P$, applying the expression from the second equation into the former equation we derive $E_2^P = (s + m_f^2 - p_{1k}^2)/(4m_\chi)$. This implies $t_1 = (k_1 - p_{1k})^2 = (k_2 - p_2)^2 = m_\chi^2 + m_f^2 - 2 \left[\frac{1}{4}(s + m_f^2 - p_{1k}^2) - m_\chi |v_\chi| \cos\theta_P \left\{ \frac{(s + m_f^2 - p_{1k}^2)^2}{16m_\chi^2} - m_f^2 \right\}^{1/2} \right]$, and $u_1 = (k_1 - p_2)^2 = (k_2 - p_{1k})^2 = m_\chi^2 + m_f^2 - 2 \left[\frac{1}{4}(s + m_f^2 - p_{1k}^2) + m_\chi |v_\chi| \cos\theta_P \left\{ \frac{(s + m_f^2 - p_{1k}^2)^2}{16m_\chi^2} - m_f^2 \right\}^{1/2} \right]$. We have used $p_2 \cdot p_{1k} = 1/2(s - m_f^2 - p_{1k}^2)$, $k_1 \cdot k_2 = s/2 - m_\chi^2$, $2k \cdot p_1 = p_{1k}^2 - k^2 - m_f^2$. Since $p_2^2 = (p_2^0)^2 - (E_2^P)^2 + m_f^2 = m_f^2$, we have $[p_2^0]_P = E_2^P = (s + m_f^2 - p_{1k}^2)/4m_\chi$. Similarly, from $P^0 = p_{1k}^0 + p_2^0$, we derive $[p_{1k}^0]_P = \sqrt{s}/2 - m_f^2/2\sqrt{s} + p_{1k}^2/2\sqrt{s}$.

By definition, we have $[p_{1k}^0]_{p_{1k}} = E_1^{p_{1k}}$, $k^2 = (k^0)^2 - (E_1^{p_{1k}})^2 + m_f^2$, and $2p_1 \cdot k = p_{1k}^2 - m_f^2 - k^2 = 2E_1^{p_{1k}} k^0 + 2((E_1^{p_{1k}})^2 - m_f^2)$. We can solve for k^0 and $E_1^{p_{1k}}$ from the latter two equations to derive $[k^0]_{p_{1k}} = \frac{p_{1k}^2 - m_f^2 + k^2}{2\sqrt{p_{1k}^2}}$ and $[E_1^{p_{1k}}]_{p_{1k}} = \frac{p_{1k}^2 + m_f^2 - k^2}{2\sqrt{p_{1k}^2}}$. By definition, we have $[p_{1k}^\mu]_{p_{1k}} = (\sqrt{p_{1k}^2}, 0, 0, 0)$, so that the Lorentz factor for

boost from the p_{1k} rest frame to the rest frame of P is $\gamma = \frac{[p_{1k}^0]_P}{[p_{1k}^0]_{p_{1k}}} = \frac{\sqrt{s}/2 - m_f^2/2\sqrt{s} + p_{1k}^2/2\sqrt{s}}{\sqrt{p_{1k}^2}}$. This implies $[p_{1k}^3]_P = \sqrt{\frac{s}{4} + \frac{m_f^4}{4s} + \frac{p_{1k}^4}{4s} - \frac{m_f^2}{2} - \frac{p_{1k}^2}{2} - \frac{m_f^2 p_{1k}^2}{2s}}$.

In the p_{1k} rest frame, conservation of momentum implies $|\mathbf{p}_1]_{p_{1k}}| = |\mathbf{k}]_{p_{1k}}|$, from which we obtain $[E_{p_1}]_{p_{1k}} = \sqrt{m_f^2 + ([E_\gamma]_{p_{1k}})^2}$, where $[E_{p_1}]_{p_{1k}}$ and $[E_\gamma]_{p_{1k}}$ denote the energy of the particle with 4-momentum p_1 and the photon in the rest frame of p_{1k} respectively. Similarly conservation of energy in the same reference frame implies $p_{1k}^0 = p_1^0 + k^0 = \sqrt{p_{1k}^2}$, from which we obtain $E_\gamma^{p_{1k}} = \frac{p_{1k}^2 - m_f^2}{2\sqrt{p_{1k}^2}}$. From all these expression, we obtain the energy of the photon in the laboratory frame as

$$\begin{aligned} [E_\gamma]_P &= \gamma[E_\gamma]_{p_{1k}} + \gamma\beta[p_\gamma^z]_{p_{1k}} \\ &= \frac{(p_{1k}^2 - m_f^2) \left(\frac{\sqrt{s}}{2} - \frac{m_f^2}{2\sqrt{s}} + \frac{p_{1k}^2}{2\sqrt{s}} \right)}{2q_1^2} \\ &\quad - \frac{(p_{1k}^2 - m_f^2) \left(\frac{s}{4} + \frac{m_f^4}{4s} + \frac{p_{1k}^4}{4s} - \frac{m_f^2}{2} - \frac{p_{1k}^2}{2} - \frac{m_f^2 p_{1k}^2}{2s} \right)}{2p_{1k}^2} \\ &\quad \times \cos\theta_q. \end{aligned} \quad (19)$$

Taking $\cos\theta_q = \pm 1$, we obtain the extremum value of p_{1k}^2 . We derive the differential cross section w.r.t. the photon energy in the laboratory frame as

$$\begin{aligned} \left[\frac{d(\sigma v)}{dE_\gamma} \right]_{\text{lab.}} &= \frac{1}{2s} \frac{1}{4} \int \sum |\mathcal{M}|^2 \frac{d(p_{1k}^2)}{2^6 (2\pi)^4} \frac{d\phi d\cos\theta_P}{s} \\ &\quad \times (-2) \sqrt{4s}. \end{aligned} \quad (20)$$

Similarly, in the center of momentum frame of p_{1k} , we obtain $[E_{p_1}]_{p_{1k}} = (p_{1k}^2 + m_f^2)/2\sqrt{p_{1k}^2}$. The energy of fermion f is

$$\begin{aligned} [E_{p_1}]_P &= \gamma[E_\gamma]_{p_{1k}} + \gamma\beta[p_\gamma^z]_{p_{1k}} \\ &= \frac{(p_{1k}^2 + m_f^2) \left(\frac{\sqrt{s}}{2} - \frac{m_f^2}{2\sqrt{s}} + \frac{p_{1k}^2}{2\sqrt{s}} \right)}{2q_1^2} \\ &\quad + \frac{(p_{1k}^2 - m_f^2) \left(\frac{s}{4} + \frac{m_f^4}{4s} + \frac{p_{1k}^4}{4s} - \frac{m_f^2}{2} - \frac{p_{1k}^2}{2} - \frac{m_f^2 p_{1k}^2}{2s} \right)}{2p_{1k}^2} \\ &\quad \times \cos\theta_q. \end{aligned} \quad (21)$$

The differential cross section w.r.t. the fermion energy in the laboratory frame as

$$\begin{aligned} \left[\frac{d(\sigma v)}{dE_{p_1}} \right]_{\text{lab.}} &= \frac{1}{2s} \frac{1}{4} \int \sum |\mathcal{M}|^2 \frac{d(p_{1k}^2)}{2^6 (2\pi)^4} \frac{d\phi d\cos\theta_P}{s} \\ &\quad \times 2\sqrt{4s}. \end{aligned} \quad (22)$$

We can obtain the photon spectrum from the fermion as follows [71]. Let us write the energy distribution of fermion per annihilation as $\frac{dN_f}{d\gamma_f} = \frac{1}{v\sigma} \frac{dv\sigma}{d\gamma_f}$, where $\gamma_f = \frac{E_f}{m_f}$. Using Pythia [73], we obtain the spectrum

of photons from fermion decay at rest $\left(\frac{dN_\gamma}{dE} \right)_{\text{at rest}}$ [73]. The energy distribution of photons from fermion arising from dark matter annihilation is

$$\begin{aligned} \left[\frac{dN}{dE} \right]_{\text{lab.}} &= \int_{-1}^{+1} \frac{d\cos\theta'}{2} \int d\gamma_f \frac{dN_f}{d\gamma_f} \int dE' \left(\frac{dN_\gamma}{dE'} \right)_{\text{at rest}} \\ &\quad \times \delta(E - [\gamma_f E' + \beta_f \gamma_f p' \cos\theta']). \end{aligned} \quad (23)$$

Integrating over $\cos\theta'$, we obtain

$$\begin{aligned} \left[\frac{dN}{dE} \right]_{\text{lab.}} &= \frac{1}{2} \int_{\gamma_f=1}^{\gamma_{\text{max}}} \frac{d\gamma_f}{\sqrt{\gamma_f^2 - 1}} \frac{dN_f}{d\gamma_f} \\ &\quad \times \int_{E'_-}^{E'_+} \frac{dE'}{E'} \left(\frac{dN_\gamma}{dE'} \right)_{\text{at rest}}, \end{aligned} \quad (24)$$

where $E'_{pm} = \gamma E \pm \gamma\beta p$ and $\gamma_{\text{max}} = m_\chi/m_f$. We multiply by an extra factor of 1/2 as $(dN_\gamma/dE)_{\text{at rest}}$ is due to a fermion and anti-fermion pair at rest from Pythia and $[dN/dE]_{\text{lab.}}$ is due to one fermion. The spectrum of photons from anti-fermion is identical.

C. Numerical Results

1. Behaviour of the 2-body and 3-body cross section

In the current section we elaborate on our analytical results. We plot the ratio of the 3-body cross section to the 2-body cross section in Fig. 2. In the top panel of Fig. 2, we show the ratio as a function of the dark matter mass for a fixed $\Lambda = 5$ TeV. We set $d_6 = 1$, $d_8 = 1$, and $v = 10^{-3}c$ for this plot. We vary the dark matter mass from $m_\chi = 1$ TeV to 5 TeV. The ratio is $\lesssim 1$ for dark matter masses less than 1 TeV. The 3-body cross section is proportional to Λ^{-8} and the 2-body cross section is proportional to Λ^{-4} , and at small dark matter masses, the enhancement of the bremsstrahlung is not able to overcome the suppression due to the 4 additional powers of Λ . As the dark matter mass increases beyond 1 TeV, the ratio starts increasing beyond 1. The ratio of the 3-body cross section to the 2-body cross section is ~ 10 for $m_\chi \sim 1.5$ TeV. For this particular choice of Λ , the cross section peaks at around $m_\chi \approx 3.5$ TeV. The maximum value of the ratio of the 3-body cross section to the 2-body cross section is $\sim 2 \times 10^6$. The ratio decreases when dark matter mass is greater than 3.5 TeV. It is interesting to note that the ratio is asymmetric around its maximal value.

In the bottom panel of Fig. 2, we show the ratio of the 3-body cross section to the 2-body cross section as a

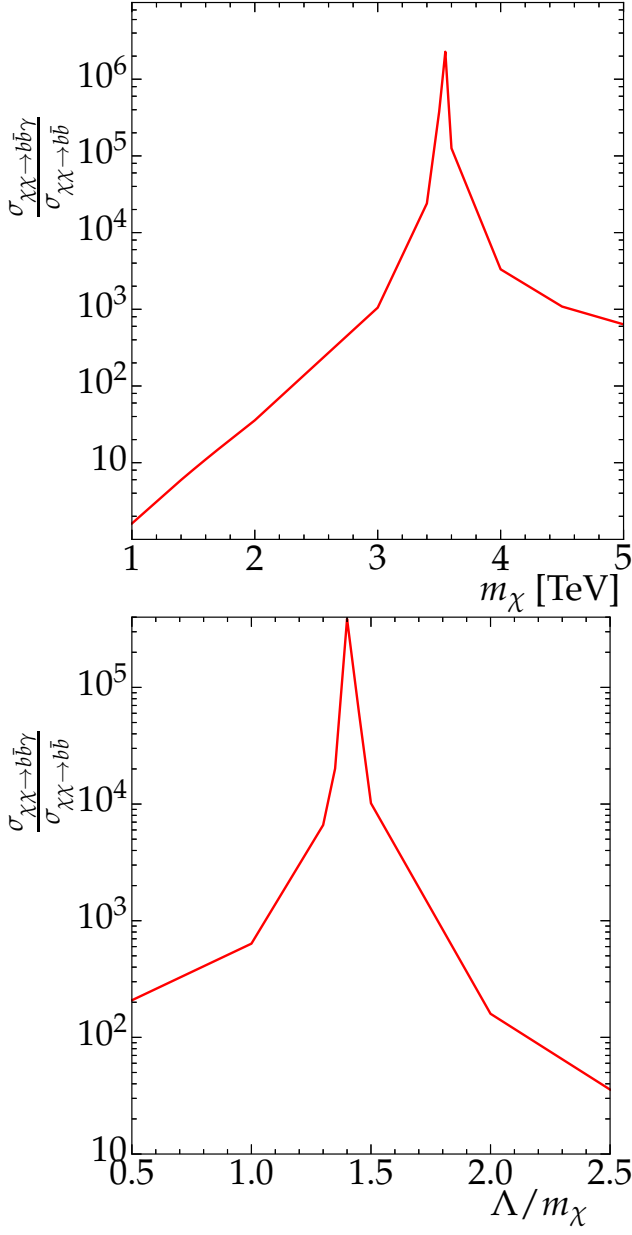


FIG. 2. Ratio of 3-body cross-section ($\chi\chi \rightarrow b\bar{b}\gamma$) to the 2-body cross-section ($\chi\chi \rightarrow b\bar{b}$) as a function of the dark matter mass, m_χ (top), and as a function of the ratio of the EFT scale to the dark matter mass, Λ/m_χ (bottom). In the top panel we set $\Lambda = 5$ TeV and in the bottom panel we set $m_\chi = 2$ TeV. We set $d_6 = 1$, $d_8 = 1$, $v = 10^{-3}c$ for both these plots.

function of the ratio Λ/m_χ . We take $m_\chi = 2$ TeV and vary Λ/m_χ from 0.5 to 2.5 for this plot. The ratio peaks at $\Lambda/m_\chi = 1.4$. It is interesting to note that the ratio of the 3-body cross section to the 2-body cross-section peaks at $\Lambda/m_\chi = 1.4$ for both the cases that we consider. This can be analytically understood as follows: for the 2-body cross section there is a cancellation between the dimension 8 term and the dimension 6 term. Setting $d_6 = d_8 = 1$ and assuming $\Lambda \gg m_f$, we find from Eq. (10) that the 2-body cross section is minimized when $\Lambda/m_\chi \approx 1.4$.

A non-zero value for the 2-body cross section comes from additional corrections to this ratio which depends on the fermion mass that we consider.

In Fig. 3 the blue (brown) dash-dotted (dashed) line corresponds to the 3-body (2-body) cross-section. The parameters assumed in the top and bottom panel of this plot are the same as that of the top and bottom panel of Fig. 2 respectively. The cancellation between

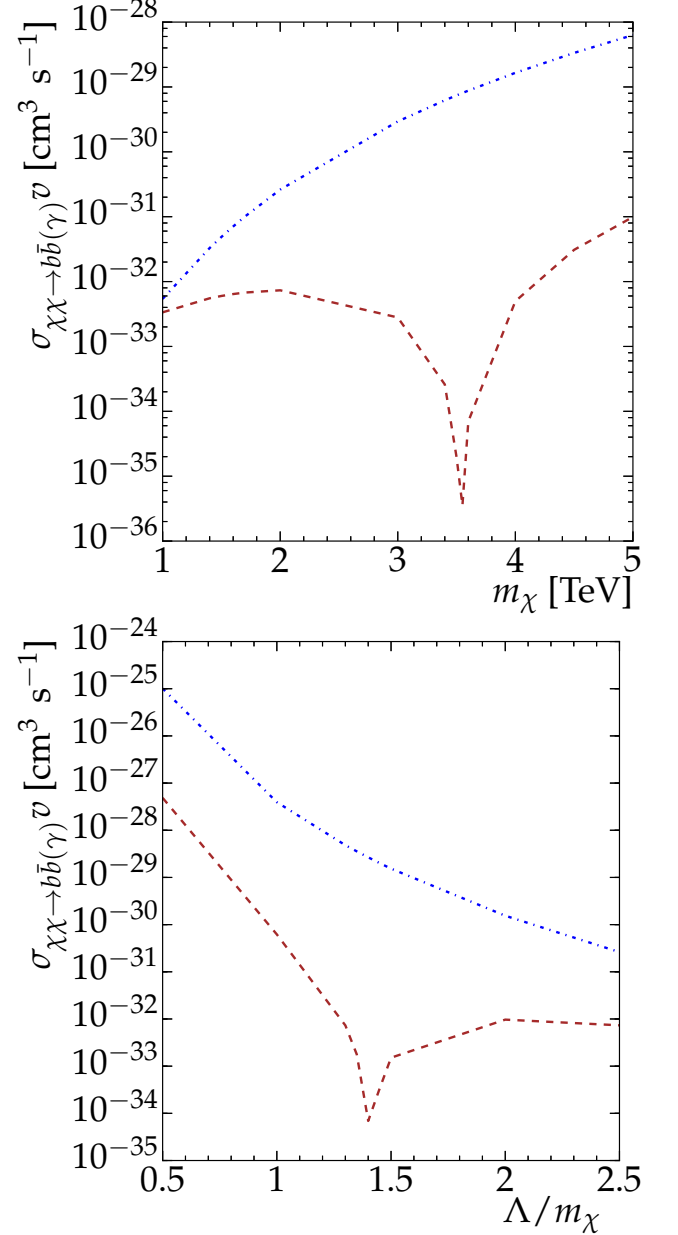


FIG. 3. Cross-sections for the 3-body and 2-body process as a function of the dark matter mass, m_χ (top), and as a function of the ratio Λ/m_χ (bottom). In the top panel we set $\Lambda = 5$ TeV and in the bottom panel $m_\chi = 2$ TeV. The 3-body cross section is denoted by the blue dash-dotted line and the brown dashed line denotes the 2-body cross section. We set $d_6 = 1$, $d_8 = 1$, and $v = 10^{-3}c$ for both of these plots.

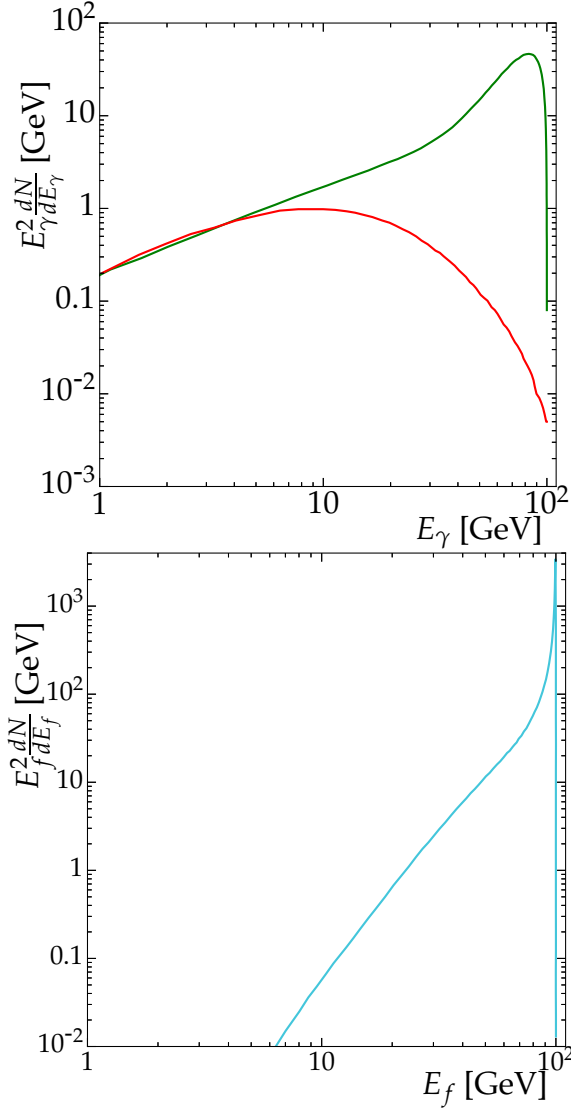


FIG. 4. Energy spectrum of the prompt γ -ray (top) and the bottom-quark (bottom) for the process $\chi\chi \rightarrow b\bar{b}\gamma$. The continuous red-line on the top panel denotes the γ -ray spectrum from the b -quark decay. In these plots we set $d_6 = 1$, $d_8 = 1$, $v = 10^{-3}c$, $m_\chi = 100$ GeV, and $\Lambda = 1$ TeV.

the dimension 6 term and the dimension 8 term for the 2-body annihilation channel $\chi\chi \rightarrow b\bar{b}$ is clearly seen in this plot. The top panel shows that the cross section for the 3-body annihilation channel $\chi\chi \rightarrow b\bar{b}\gamma$ monotonically increases with increasing m_χ for a fixed Λ . The bottom panel shows that the cross section for the 3-body annihilation channel $\chi\chi \rightarrow b\bar{b}\gamma$ monotonically decreases for an increasing Λ when m_χ is kept fixed.

2. Energy spectrum of the final state photons and fermions

The energy spectrum of the fermion emitted at the production vertex can be obtained from the cross-section

of the process as follows

$$\frac{dN_f}{dE_f} = \frac{1}{v\sigma} \frac{d(v\sigma)}{dE_f}, \quad (25)$$

where E_f denotes the final energy of the fermion. The energy spectrum of the photons radiated from the final state fermion is obtained using Pythia [73]. The events are simulated at the center-of-mass energy $2m_\chi$. In Fig. 4 we show the energy spectrum of the prompt photon (top panel) and the final state fermion (bottom panel) for the annihilation channel $\chi\chi \rightarrow b\bar{b}\gamma$. The continuous green line in the top panel of Fig. 4 depicts the spectrum of the prompt photon, while the continuous red line shows the spectrum of the photons emitted from the final state bottom quark. In Fig. 4, we use the dark matter particle mass as $m_\chi = 100$ GeV, $\Lambda = 1$ TeV, $d_6 = d_8 = 1$, and $v = 10^{-3}c$.

We show the energy spectrum as $E^2 dN/dE$ in a log-log plot. The area under the curve gives the total energy carried by the particle. From the top panel in Fig. 4, we see that the prompt photon carries more energy than the energy carried by the photons coming from b -quark hadronisation and decay. This can be understood intuitively as the b -quark hadronisation and decay produces charged anti-particles and neutrinos which also carry away a substantial portion of the energy of the b -quark.

The smooth spectrum of the photons resulting from b -quark hadronisation and decay imply that it will be difficult to distinguish this spectrum from the spectrum of conventional astrophysical sources. On the other hand, the spectrum of the prompt photons is quite unlike anything produced in astrophysics, and it will be easier to distinguish this spectrum from the background produced by conventional astrophysical sources.

The bottom panel of Fig. 4 shows the energy distribution of the final state b -quark. We see from Fig. 4 that most of the energy is coming from the region around m_χ . While there is a sharp feature in the bottom-quark energy distribution, in practice it will not be visible as the bottom-quark hadronises and decays after its production. The spectrum of the decay products shows a smooth feature which is hard to discriminate from the spectrum of conventional astrophysical sources.

We do not show the spectrum of the electrons, positrons, protons, anti-protons, neutrinos and anti-neutrinos arising from b -quark hadronisation and decay. They can be obtained in a similar way as we obtained the spectrum of the photons from b -quark hadronisation and decay.

3. Constraints on Λ

Fig. 5 compares the constraint on the effective operator scale Λ from various different experiments as a function of the dark matter mass. For the indirect detection

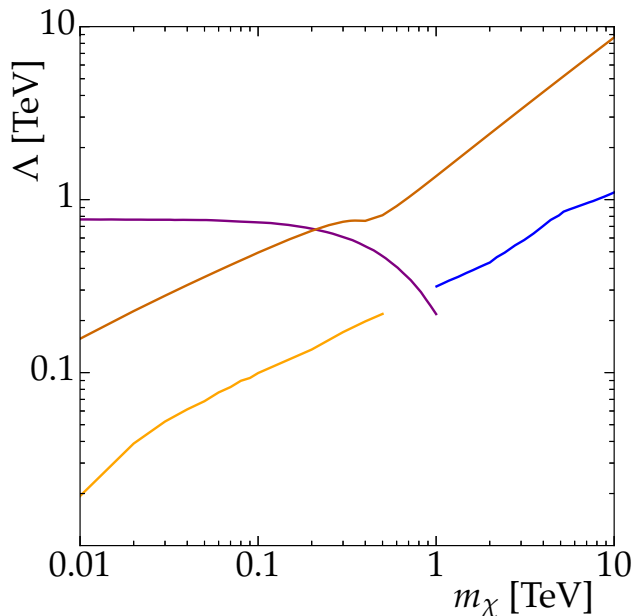


FIG. 5. Lower bound on the effective operator scale Λ as a function of the dark matter mass m_χ from indirect detection experiments and collider experiments. The orange and blue lines denote the bound derived from γ -ray fluxes from Fermi-LAT [14], and MAGIC [74] experiment. The dark matter relic density, as measured by the Planck collaboration [75], is satisfied when Λ lies on the brown line. The purple line is the bound coming from mono- γ searches at the collider [40] for the operator $\mathcal{L}_{d=6}$. If Λ lies below the brown line then dark matter particles make up a fraction of the relic density using the 3-body annihilation channel. We set $d_6 = d_8 = 1$.

experiments, we will only consider constraints from gamma-ray experiments. Constraints from anti-protons, and positrons are competitive, however, they are subject to additional uncertainties due to our less than precise knowledge on transport and energy loss properties of the charged anti-particles. Constraints from neutrinos are weaker compared to the upper limits from gamma-ray, and charged particle experiments.

Gamma-ray experiments typically only present the constraints on 2-body annihilation processes. However, Fermi-LAT and MAGIC experiment presented the bin-by-bin upper limit from their experiments in Refs. [14] and [74] respectively. We compare these bin-by-bin upper limits with our derived total gamma-ray spectrum to derive the constraints on Λ .

The orange and blue line is the bound obtained from γ -ray fluxes from Fermi-LAT [14] and MAGIC [74] experiment respectively. For a given m_χ the region below these lines will have lower Λ and will produce much more γ -ray fluxes than the observed one, thus they are disfavored. While the relic density is inversely proportional to the cross-section, and thus it is proportional to some power of the effective operator scale. So, the region above the brown line will have large relic abundance of the dark matter which is in

contradiction to the Planck collaboration [75] results. For a given m_χ , on the brown line one satisfies the relic density, whereas in the region below the brown line relic density is lower than the Planck measurement and to address the Planck measured relic abundance, here one needs to invoke the idea of multi-component dark matter or non-thermal production of the the dark matter.

To compare this limit as obtained from the indirect detection experiments, with the direct detection experiments like the LHC, we show the purple line in Fig. 5. The purple line denotes the lower bound on the Λ from the mono- γ searches at the LHC for $d = 6$ operator only (see Eq.(1)) [40]. Thus the region below the purple line is disfavored from the mono- γ searches at the collider. This is in contradiction to the relic density limit until the mono- γ search limit crosses the relic density limit at m_χ 200 GeV or so. Hence we can see that above $m_\chi > 200$ GeV there is a common allowed region which is safe from both indirect and collider searches.

The validity of the effective field theory paradigm depends on the details of the coupling constants [44, 76]. In general, one can take it to be $m_\chi \lesssim \Lambda$, where the inequality can change by a factor of few depending on the underlying UV-complete model. The value of Λ that we derive when we assume that the 3-body annihilation process makes up the full relic density is consistent with the regime of validity. The indirect detection limits that we present are in a regime in which one can question the validity of the effective field theory. This implies that if one sees a signal in the very near future in indirect detection experiments, then the effective field theory approach will not be a good description of the signal for the model that we have considered.

III. CONCLUSION

In this paper we have considered the annihilation of two Majorana dark matter particle into SM light fermions. The annihilation of self-conjugate dark matter particles to a pair of fermions is helicity suppressed, but a photon radiation lifts up this suppression. This opens up the s -wave channel irrespective of the dark matter relative velocity. We addressed this phenomena from an effective operator point of view, by using a dimension 8 operator.

We have calculated the cross-section of the dark matter annihilation to light fermions and a photon in the presence of dimension 6 and dimension 8 operator. We have found that in spite of the higher dimensionality of the dimension 8 operator, it does not suffer from any suppression due to dark matter relative velocity. We have showed that the contribution to the annihilation cross-section from dimension 8 operator to the dimension 6 operator is always larger at all dark matter mass scales $\gtrsim 1$ TeV. Also, we have shown that there is a cancellation in the 2-body cross-section between the pure dimension

6 and dimension 8 to their interference terms for $d_6 = 1$, $d_8 = 1$, and $\Lambda/m_\chi \simeq 1.4$.

We have shown the photons from VIB receives dominant contribution from prompt decay rather than the decay of the final state fermions. We then calculated the bounds on the EFT scale as a function of the dark matter mass from various indirect detection experiments like, γ -ray fluxes measured at the Fermi-LAT and MAGIC experiment and the dark matter relic density measured by the Planck collaboration. We have compared these bounds with the mono- γ searches at the collider. We have found that for low dark matter mass ($\lesssim 200$ GeV) the mono- γ searches supersedes the bound coming from the indirect detection experiments. Whereas above 200 GeV dark matter mass the relic

density constraint provides the strongest bound on the EFT scale.

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